

A comparative study of Jet-quenching Schemes

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Abstract. The four major approximation schemes devised to study the modification of jets in dense matter are outlined. The comparisons are restricted to basic assumptions and approximations made in each case and the calculation methodology used. Emergent underlying similarities between apparently disparate methods brought about by the approximation schemes are exposed. Parameterizations of the medium in each scheme are discussed in terms of the transport coefficient \hat{q} . Discrepancies between the estimates obtained from the four schemes are discussed. Recent developments in the basic theory and phenomenology of energy loss are highlighted.

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1. Introduction

The considerable modification of the spectrum of high transverse momentum (p_T) hadrons produced in ultra-high energy heavy-ion collisions has now been established by experiments at the Relativistic Heavy-Ion Collider (RHIC). The number of such hadrons, irrespective of flavour, with $p_T \geq 7\text{GeV}$ is reduced by almost a factor of 5 in central Au - Au collisions compared to that expected from elementary nucleon nucleon encounters enhanced by the number of expected binary collisions [1, 2]. This suppression, referred to as *jet-quenching*, caused by the energy loss of very energetic partons produced in the few hard scatterings that occur early in the collision, represents one of the major theoretical predictions which have been substantiated by experiment [3, 4]. The basic underlying mechanism, that of induced gluon radiation from the hard parton traversing a coloured environment, has by now been expounded upon in multiple reviews [5, 6] and there exists a certain consensus regarding the physics involved.

In the last several years, dramatic progress has been made both in experiment and theory. While measurements on single inclusive observables have been extended to wider ranges in p_T and over a variety of systems, there has arisen a plethora of multiparticle jet-like correlation observables, photon-jet, jet-medium and heavy flavour observables [7]‡. On the theory side, calculations have evolved through multiple levels of complexity and have also been expanded in scope to address the wide variety of new observables. In the following, the status of energy loss schemes in the light flavour sector are reviewed. New observables capable of offering deeper insight into the structure of

‡ See also Ref. [8] and references therein

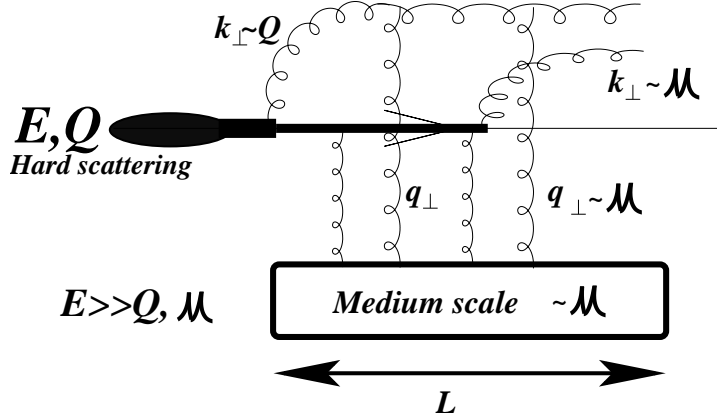


Figure 1. A schematic picture of the various scales involved in the modification of jets in dense matter.

the dense matter produced and differentiating between the various schemes, as well as new developments in the theory of jet modification and correlations are outlined.

2. The four schemes

The majority of current approaches to the energy loss of light partons may be divided into four major schemes often referred to by the names of the original authors. All schemes utilize a factorized approach where the final cross section to produce a hadron h with transverse momentum p_T (rapidity between y and $y + dy$) may be expressed as a convolution of initial nuclear structure functions $[G_a^A(x_a), G_b^B(x_b)$, initial state nuclear effects such as shadowing and Cronin effect are understood to be included] to produce partons with momentum fractions x_a, x_b , a hard partonic cross section to produce a high transverse momentum parton c with a transverse momentum \hat{p} and a medium modified fragmentation function for the final hadron $[\tilde{D}_c^h(z)]$,

$$\frac{d^2\sigma^h}{dyd^2p_T} = \frac{1}{\pi} \int dx_a \int dx_b G_a^A(x_a) G_b^B(x_b) \frac{d\sigma_{ab \rightarrow cX}}{d\hat{t}} \frac{\tilde{D}_c^h(z)}{z}. \quad (1)$$

In the vicinity of mid-rapidity, $z = p_T/\hat{p}$ and $\hat{t} = (\hat{p} - x_a P)^2$ (P is the average incoming momentum of a nucleon in nucleus A). The entire effect of energy loss is concentrated in the calculation of the modification to the fragmentation function. The four models of energy loss are in a sense four schemes to estimate this quantity from perturbative QCD calculations. To better appreciate the approximation schemes, one may introduce a set of scales (see Fig. 1): E or p^+ , the forward energy of the jet, Q^2 , the virtuality of the initial jet-parton, μ , the momentum scale of the medium and L , its spatial extent. Most of the differences between the various schemes may be reduced to the different relations between these various scales assumed by each scheme as well as by how each scheme treats or approximates the structure of the medium. In all schemes, the forward energy of the jet far exceeds the medium scale, $E \gg \mu$. The schemes are presented from one extreme of the approximation set (higher twist approach) to the opposite extreme (finite

temperature approach), similarities in intermediate steps of the calculation will not be repeated.

2.1. Higher Twist

The origin of the higher twist (HT) approximation scheme lies in the calculations of medium enhanced higher twist corrections to the total cross section in Deep-Inelastic Scattering (DIS) off large nuclei [9]. The essential idea lies in the inclusion of power corrections to the leading twist cross sections, which, though suppressed by powers of the hard scale Q^2 , are enhanced by the length of the medium. This technology of identifying and isolating power corrections is used to compute the single hadron inclusive cross-section.

One assumes the hierarchy of scales $E \gg Q \gg \mu$ and applies this to the computation of multiple Feynman diagrams such as the one in the left panel of Fig. 2, this diagram represents the process of a hard virtual quark produced in a hard collision, which then radiates a gluon and then scatters off a soft medium gluon with transverse momentum $q_\perp \sim \mu$ prior to exiting the medium and fragmenting into hadrons. At a given order, there exist various other contributions which involve scattering of the initial quark off the soft gluon field prior to radiation as well as scattering of the radiated gluon itself. All such contributions are combined coherently to calculate the modification to the fragmentation function directly.

The hierarchy of scales allows one to use the collinear approximation to factorize the fragmentation function and its modification from the hard scattering cross section. Thus, even though such a modified fragmentation function is derived in DIS, it may be generalized to the kinematics of a heavy-ion collision. Diagrams where the outgoing parton scatters off the medium gluons, such as those in Fig. 2, produce a medium dependent additive contribution to the vacuum fragmentation function, which may be expressed as,

$$\Delta D_i(z, Q^2) = \int_0^{Q^2} \frac{dk_\perp^2}{k_\perp^2} \frac{\alpha_s}{2\pi} \left[\int_{z_h}^1 \frac{dx}{x} \sum_{j=q,g} \left\{ \Delta P_{i \rightarrow j}(x, x_B, x_L, k_\perp^2) D_j^h \left(\frac{z_h}{x} \right) \right\} \right]. \quad (2)$$

In the above equation, $\Delta P_{i \rightarrow j}$ represents the medium modified splitting function of parton i into j where a momentum fraction x is left in parton j . The factor, $x_L = k_\perp^2 / (2P^- p^+ x(1-x))$ §, where the radiated gluon or quark carries away a transverse momentum k_\perp , P^- is the incoming momentum of a nucleon in the nucleus and p is the momentum of the virtual photon. The medium modified splitting functions may be expressed as a product of the vacuum splitting function $P_{i \rightarrow j}$ and a medium dependent factor,

§ Throughout these proceedings, four-vectors will often be referred to using the light cone convention where $x^\pm = (x^0 \pm x^3)/\sqrt{2}$. For the higher-twist scheme, often, $x^+ = (x^0 + x^3)/2$ and $x^- = x^0 - x^3$.

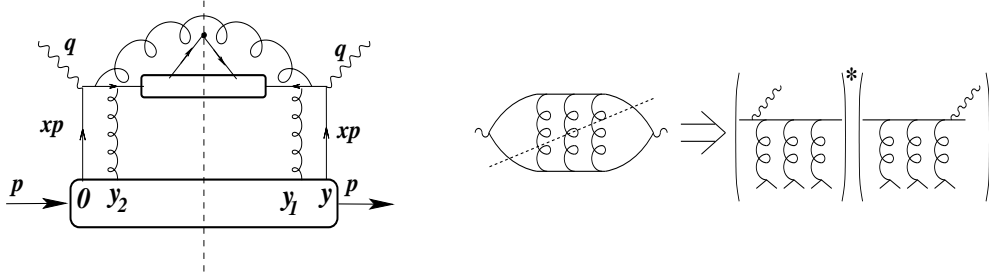


Figure 2. Left panel: a higher twist contribution to the modification of the fragmentation function in medium. Right panel: a typical cut diagram in the AMY formalism.

$$\Delta \hat{P}_{i \rightarrow j} = P_{i \rightarrow j} \frac{C_A 2\pi\alpha_s T_{qq}^A(x_a, x_L)}{(k_\perp^2 + \langle q_\perp^2 \rangle) N_c f_q^A(x_a)}. \quad (3)$$

Where, C_A, N_c represent the adjoint Casimir and the number of colours. The mean transverse momentum of the soft gluons is represented by the factor $\langle q_\perp^2 \rangle$. The term T_{qq}^A represents the quark gluon correlation in the nuclear medium, and depends on the four point correlator,

$$\langle P | \bar{\psi}(0) \gamma^+ F_\sigma^+(y_2^-) F^{+\sigma}(y_1^-) \psi(y^-) | P \rangle \sim C \langle p_1 | \bar{\psi}(0) \gamma^+ \psi(y^-) | p_1 \rangle \langle p_2 | F_\sigma^+(y_2^-) F^{+\sigma}(y_1^-) | p_2 \rangle. \quad (4)$$

Where, $F_\sigma^+(y_2^-)$ and $F^{+\sigma}(y_1^-)$ represent gluon field operators at the locations y_1^-, y_2^- and $\psi(y)$ represents the quark field operator. The above correlation function cannot be calculated from first principles without making assumptions regarding the structure of the medium. The only assumption made is that the colour correlation length is small. As a result, one may factorize the four point function into two separate structure functions, one for the original parton produced in the hard scattering [this is a quark in Eq. (4)] and one for the soft gluon off which the parton scatters in the final state.

The entire phenomenology of the medium is incorporated as a model for the expectation of the second set of operators in Eq. (4). The product of its absolute magnitude as well as the correlation constant C is set by fitting to one data point. Unlike the remaining formalisms, the H-T approach is set up to directly calculate the medium modified fragmentation function and as a result the final spectrum. The determined constant, C , may be used to calculate the average energy loss encountered by a jet. Another advantage of this approach is the straightforward generalization to multiparticle correlations [11] and their modification in the medium.

2.2. Opacity expansion: Reaction Operator Approach

Unlike the higher twist or the finite temperature approaches, this formalism, often referred to as the Gyulassy-Levai-Vitev (GLV) scheme [12], and other opacity expansion schemes were constructed primarily to deal with the problem of energy loss in dense deconfined matter. The GLV scheme assumes the medium to be composed of heavy

almost static colour scattering centers which are well separated in the sense that the mean free path of a jet $\lambda \gg 1/\mu$ the colour screening length of the medium [13]. The opacity of the medium \bar{n} quantifies the number of scattering centers seen by a jet as it passes through the medium, *i.e.*, $\bar{n} = L/\lambda$, where L is the thickness of the medium. The opacity or gluon number density is the quantity used to model the presence of the medium.

A hard jet is produced locally in such a plasma with a large forward energy $p^+ \gg \mu$ and almost immediately begins to shower soft gluons. The typical transverse momentum of the radiated gluons is similar in order of magnitude to the transverse momentum imparted from the medium *i.e.*, $E \gg Q \sim \mu$. The colour centers are assumed to produce a screened Yukawa potential. At first order in opacity, a jet scatters off one such potential and picks up a transverse momentum \vec{q}_\perp ; in the process it will radiate a gluon with momentum $k \equiv (xp^+, \frac{k_\perp^2}{xp^+}, \vec{k}_\perp)$. The scattering may happen before or after the radiation. Squaring such contributions and including interference terms between vacuum radiation and double scattering, leads, in the limit of $x \rightarrow 0$ (and ignoring spin effects), to the soft gluon differential emission distribution at first order in opacity [12],

$$x \frac{dN}{dx dk_\perp^2} = x \frac{dN}{dx dk_\perp^2} \frac{L}{\lambda_g} \int_0^{q_{max}^2} d^2 q_\perp \frac{\mu_{eff}^2}{\pi(q_\perp^2 + \mu^2)^2} \frac{2k_\perp \cdot q_\perp (k - q_1)^2 L^2}{16x^2 E^2 + (k - q)_\perp^2 L^2}. \quad (5)$$

In the above equation, λ_g is the mean free path of the radiated gluon. Consideration of single inclusive gluon emission from multiple scattering requires the use of a recursive diagrammatic procedure [14]. The inclusion of such diagrams allows for the computation of gluon distributions to finite order ($n \geq 1$) in opacity.

Due to the soft limit *i.e.*, $x \rightarrow 0$ used, multiple gluon emissions are required for a substantial amount of energy loss. Each such emission at a given opacity is assumed independent and a probabilistic scheme is set up, wherein, the jet loses an energy fraction ϵ in n tries with a Poisson distribution [15],

$$P_n(\epsilon, P^+) = \frac{e^{-\langle N_g \rangle}}{n!} \Pi_{i=1}^n \left[\int dx_i \frac{dN_g}{dx_i} \right] \delta(\epsilon - \sum_{i=1}^n x_i), \quad (6)$$

where, $\langle N_g \rangle$ is the mean number of gluons radiated per coherent interaction set. Summing over n gives the probability $P(\epsilon)$ for an incident jet to lose a momentum fraction ϵ due to its passage through the medium. This is then used to model a medium modified fragmentation function, by shifting the energy fraction available to produce a hadron (as well as accounting for the phase space available after energy loss),

$$\tilde{D}(z, Q^2) = \int_0^1 d\epsilon P(\epsilon) \frac{D(\frac{z}{1-\epsilon}, Q^2)}{1-\epsilon}. \quad (7)$$

The above, modified fragmentation function is then used in a factorized formalism as in Eq. (1) to calculate the final hadronic spectrum.

2.3. Opacity expansion: Path Integral Approach

The path integral approach for the energy loss of a hard jet propagating in a coloured medium was first introduced in Ref. [16]. It was later demonstrated to be equivalent to the well known BDMPS approach [17]. The current, most widespread, variant of this approach developed by numerous authors is often referred to as the Armesto-Salgado-Wiedemann (ASW) approach. In this scheme, one incorporates the effect of multiple scattering of the incoming and outgoing partons in terms of a path integral over a path ordered Wilson line [18].

Similar to the GLV approach, this formalism also assumes a model for the medium as an assembly of Debye screened heavy scattering centers. A hard, almost on shell, parton traversing such a medium will engender multiple transverse scatterings of order $\mu \ll p^+$. It will in the process split into an outgoing parton and a radiated gluon which will also scatter multiply in the medium. The radiation, being brought about by the multiple scattering has a transverse momentum $k_\perp \sim \mu$ (similar to the GLV and different from the HT approach). The propagation of the incoming (outgoing) partons as well as that of the radiated gluon in this background colour field may be expressed in terms of effective Green's functions $[G(\vec{r}_\perp, z; \vec{r}'_\perp, z')]$ (for quark or gluon) which obey the obvious Dyson-Schwinger equation,

$$G(\vec{r}_\perp, z; \vec{r}'_\perp, z') = G_0(\vec{r}_\perp, z; \vec{r}'_\perp, z') - i \int_z^{z'} d\zeta \int d^2\vec{x} G_0(\vec{r}_\perp, z; \vec{x}, \zeta) A_0(\vec{x}, \zeta) G(\vec{x}, \zeta; \vec{r}'_\perp, z'), \quad (8)$$

where, G_0 is the free Green's function and A_0 represents the colour potential of the medium. The solution for the above interacting Green's function involves a path ordered Wilson line which follows the potential from the location $[\vec{r}_\perp(z'), z']$ to $[\vec{r}_\perp(z), z]$. Expanding the expression for the radiation cross section to order A_0^{2n} corresponds to an expansion up to n^{th} order in opacity.

Taking the high energy limit and the soft radiation approximation ($x \ll 1$), one focuses on isolating the leading behaviour in x that arises from the large number of interference diagrams at a given order of opacity. As a result of the approximations made, one recovers the BDMPS condition that the leading behaviour in x is contained solely in gluon re-scattering diagrams. This results in the expression for the inclusive energy distribution for gluon radiation off an in-medium produced parton as [19],

$$x \frac{dI}{dx} = \frac{\alpha_s C_R}{(2\pi)^2 x^2} 2\text{Re} \int_{\zeta_0}^{\infty} dy_l \int_{y_l}^{\infty} d\bar{y}_l \int_0^{\chi x p^+} d\vec{u} \int d\vec{k} e^{-i\vec{k} \cdot \vec{u} - \frac{1}{2} \int d\zeta n(\zeta) \sigma(\vec{u})} \frac{\partial^2}{\partial y \partial u} \int_{\vec{y}=0=\vec{r}(y_l)}^{\vec{u}=\vec{r}(\bar{y})} \mathcal{D} r e^{i \int d\zeta \frac{x p^+}{2} \left(|\vec{r}|^2 - \frac{n(\zeta) \sigma(\vec{r})}{i x p^+} \right)}, \quad (9)$$

where, as always, k_\perp is the transverse momentum of the radiated gluon and $x p^+$ is its forward momentum. The vectors \vec{y} and \vec{u} represent the transverse locations of the emission of the gluon in the amplitude and the complex conjugate whereas y_l and \bar{y}_l represent the longitudinal positions. The density of scatterers in the medium at location

ζ is $n(\zeta)$ and the scattering cross section is $\sigma(r)$. In this form, the opacity is obtained as $\int n(\zeta)d\zeta$ over the extent of the medium.

Numerical implementations of this scheme have focused on two separate regimes. In one case, $\sigma(r)$ is replaced with a dipole form Cr^2 and one solves the harmonic oscillator like path integral. This corresponds to the case of multiple soft scatterings of the hard probe. In the other extreme, one expands the exponent as a series in $n\sigma$; keeping only the leading order term corresponds to the picture of gluon radiation associated with a single scattering. In this second form, the analytical results of the ASW scheme formally approach those of the GLV reaction operator expansion [19].

In either case, the gluon emission intensity distribution has been found to be rather similar, once scaled with the characteristic frequency in each case. Beyond this, the calculation of the total energy loss in the ASW scheme follows a procedure similar to that in the GLV approach, where, a probabilistic scheme to lose an energy fraction ϵ in multiple independent emissions is set up, as in Eq. (6). This is used to calculate the medium modified fragmentation function [20], as in Eq. (7).

2.4. Finite temperature field theory approach

In this scheme, often referred to as the Arnold-Moore-Yaffe (AMY) approach, the energy loss of hard jets is considered in an extended medium in equilibrium at asymptotically high temperature $T \rightarrow \infty$. Due to asymptotic freedom, the coupling constant $g \rightarrow 0$ at such high temperatures and a power counting scheme emerges from the ability to identify a hierarchy of parametrically separated scales $T \gg gT \gg g^2T$ etc. In this limit, it becomes possible to construct an effective field theory of soft modes, *i.e.*, $p \sim gT$ by summing contributions from hard loops with $p \sim T$, into effective propagators and vertices [21].

One assumes a hard on-shell parton, with energy several times that of the temperature, traversing such a medium, undergoing soft scatterings with momentum transfers $\sim gT$ off other hard partons in the medium. Such soft scatterings induce collinear radiation from the parton, with a transverse momentum of the order of gT . The formation time for such collinear radiation $\sim 1/(g^2T)$ is of the same order of magnitude as the mean free time between soft scatterings [22]. As a result, multiple scatterings of the incoming (outgoing) parton and the radiated gluon need to be considered to get the leading order gluon radiation rate. One essentially calculates the imaginary parts of infinite order ladder diagrams such as those shown in the right panel of Fig. 2; this is done by means of integral equations [23].

The imaginary parts of such ladder diagrams yield the $1 \rightarrow 2$ decay rates of a hard parton (*a*) into a radiated gluon and another parton (*b*) Γ_{bg}^a . These decay rates are then used to evolve hard quark and gluon distributions from the initial hard collisions, when they are formed, to the time when they exit the medium, by means of a Fokker-Planck like equation [24], which is written schematically as,

$$\frac{dP_a(p)}{dt} = \int dk \sum_{b,c} \left[P_b(p+k) \frac{d\Gamma_{ac}^b(p+k, p)}{dkdt} - P_a(p) \frac{d\Gamma_{bc}^a(p, k)}{dkdt} \right]. \quad (10)$$

The initial distributions are taken from a factorized hard scattering cross section as in Eq. (1). A medium modified fragmentation function is modelled by the convolution of the vacuum fragmentation functions with the hard parton distributions, at exit, to produce the final hadronic spectrum [25],

$$\tilde{D}^h(z) = \int dp_f \frac{z'}{z} \sum_a P_a(p_f; p_i) D_a^h(z'). \quad (11)$$

Where, the sum over a is the sum over all parton species. The two momentum fractions are $z = p_h/p_i$ and $z' = p_h/p_f$, where p_i and p_f are the momenta of the hard partons immediately after the hard scattering and prior to exit from the medium. The integral above depends on the path taken by the parton through the medium, which in turn depends on the location of origin of the jet and its angle with respect to the reaction plane. The model of the medium is essentially contained in the space-time profile chosen for the temperature.

3. Comparisons to data and future directions

At the energies of collision at RHIC, complete jets cannot be resolved from the large background of particles produced; energy loss is deduced from a leading particle analysis. The primary observable in this regard is the nuclear modification factor for high p_T particles as a function of p_T and as a function of the centrality of collision. All four schemes have made successful comparisons to the available data as shown in Figs. 3,4. While the qualitative descriptions of the medium differ between the various models, all such descriptions may be subjected to comparison via a single parameter: the transport coefficient \hat{q} , defined as the average squared transverse momentum imparted to the jet per unit path length traversed in the medium.

In all four schemes, this one parameter or one directly related to it is tuned to fit the data. In the ASW scheme, \hat{q} is the fit parameter. In the GLV scheme, it depends on the density of scattering centers (dN/dy , which is the fit parameter) and the screening length in the medium. In the higher twist scheme, \hat{q} depends on the gluon density of the medium, which constitutes the fit parameter. In the AMY scheme, the fit parameter is the temperature (T) and \hat{q} depends on T . It is in these derived values of \hat{q} that differences arise between the various schemes. The GLV scheme uses an effective length as a substitute for a full space-time density profile and quotes a value of $\hat{q} \leq 1\text{GeV}^2/\text{fm}$, averaged over the space-time profile of the collision. The closely related ASW scheme quotes a value of $\hat{q} \leq 1\text{GeV}^2/\text{fm}$ for calculations employing an a similar effective length of $\sim 6\text{ fm}$. However, for realistic geometries, a value of $5\text{--}15\text{GeV}^2/\text{fm}$ is deduced. The H-T scheme quotes a value of $\sim 1\text{--}2\text{GeV}^2/\text{fm}$ as the maximum value at a time $\tau = 1\text{fm}$,

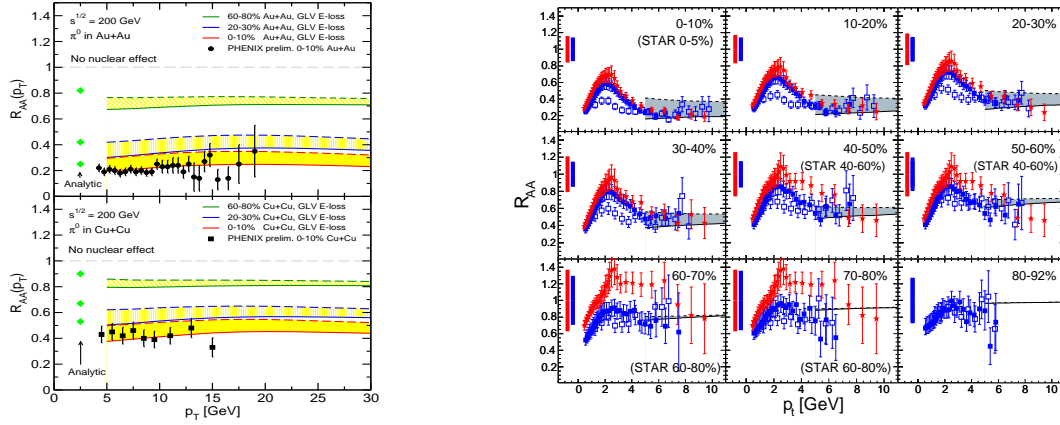


Figure 3. R_{AA} as a function of p_T and centrality in GLV (left) and ASW (right) compared to experimental data. The ASW plots used correspond to the PQM version of Ref. [26]. The GLV plots are from Ref. [27].

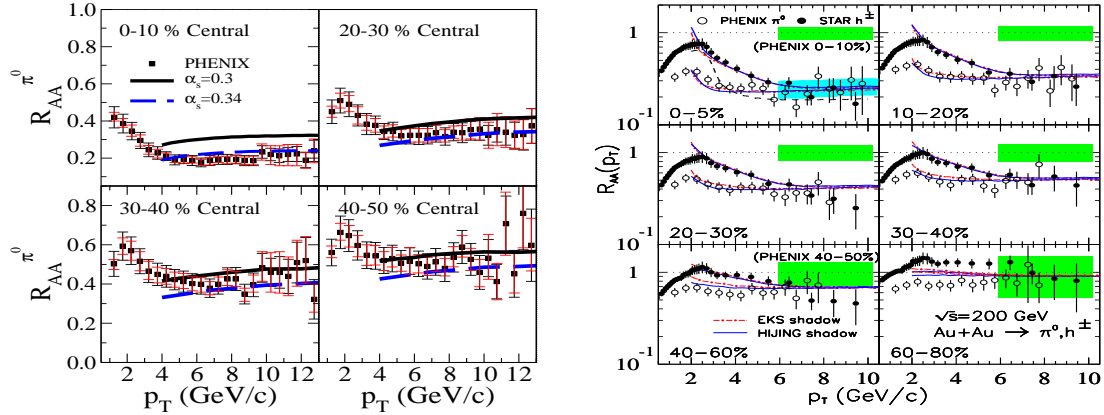


Figure 4. R_{AA} as a function of p_T and centrality in AMY [28] (left panel) and HT [29] (right panel) compared to experimental data.

while the AMY approach calculates a $\hat{q} \sim 2\text{GeV}^2/\text{fm}$ at a $T \simeq 300\text{MeV}$; in both cases, this implies a $\hat{q} \leq 1\text{GeV}^2/\text{fm}$ averaged over the life-time of the collision.

Two comments are in order: of the four schemes, the ASW and the GLV are most closely related in terms of basic physical picture and calculational scheme used whereas the HT and the AMY approaches are notably different in their origins. As a result, the difference in the estimates of \hat{q} between the ASW and GLV approaches and the similarity between the AMY and the HT schemes is remarkable. A heuristic explanation for the former has been provided by the ASW authors: in a calculation with realistic geometry, the average path length traversed by a jet was found to be $L \sim 2\text{ fm}$, three times smaller than that used in the fixed length estimates; as energy loss depends quadratically on L , the effective \hat{q} extracted was nine times larger. The second observation (regarding AMY and HT) may be understood as the result of the series of approximations made in each case. It has been demonstrated that the high temperature medium described by the HTL effective theory may also be understood in terms of a Vlasov theory of hard particles with momenta $\sim T$ moving under the influence of soft $\sim gT$ fields [30]. The jet parton is identified as similar to the hard modes $\sim T$ in terms of its propagation through

the medium. A similar picture of hard partons moving under the influence of soft fields via the effect of a colour Lorentz force, also arises in an all twist re-summation, if it is assumed that the colour correlation length in the medium is short [31].

However, the radiation spectrum as well as the hierarchy of scales used in each of the different approaches remains different. Single inclusive observables such as the R_{AA} reveal an overall effect of energy loss, where the differences between the various schemes may be masked by a renormalization of the undetectable properties of the medium. To demonstratively distinguish between the various schemes requires the use of differential probes (beyond leading hadrons) emanating from the same jet. To date, such calculations have been performed at the integrated level [11, 32]. In the left panel of Fig. 5, calculations of the variation of the integrated yield of hadrons, associated with a hard trigger hadron, as a function of centrality of the collision are plotted. The yield, naturally depends on the detected flavour content, which depends on the ability of the detector to account for decay corrections. The theoretical curves represent two extreme possibilities, where only charged pions are detected or all charged hadrons are detected including contributions from all decay corrections. The slight rise with centrality in the theoretical curves is due to the larger amount of near side energy loss in more central collisions. Yet another method to test the scheme dependence of energy loss is via single inclusive probes, subjected to a more differential analysis *e.g.*, the R_{AA} versus the reaction plane [33]. Such measurements will not only resolve the differences in medium profiles between the various schemes, but produce more detailed probes of the evolving medium as seen at the scale of the jet.

The study of energy loss in dense matter encompasses, not merely a study of a single medium property \hat{q} in terms of its effect on parton propagation, but through hard-soft correlations, can be extended to a study of how the lost energy is redistributed by the medium. Phenomenological studies of hard-soft correlations indicate that the lost energy seems to excite collective density excitations in the medium [34]. However, alternate explanations in terms of Cherenkov radiation [35] or in-medium Sudakov effects in jet radiation [36] have not been ruled out. While such a study lies outside the reach of perturbative QCD, a precise determination of the mechanism of energy loss and its distribution along the path of the jet become crucial to all model studies of energy redistribution. The requirement of unambiguous results, necessarily confines the study of energy loss mechanisms and associated measurements beyond the reach of final state bulk-medium effects such as recombination, *i.e.*, $p_T \geq 7\text{GeV}$. Such detailed studies will ultimately require higher statistics in the experimental measurements at high transverse momenta and predictions from all four schemes for the same measurement.

In spite of its long history, developments in the basic theory of jet modification continue to be made. It has recently been proposed that \hat{q} may possess a tensorial structure [37], where the scalar $\hat{q} = \delta^{ij}\hat{q}_{ij}$. An example where such a situation may arise is in the presence of large turbulent colour fields, which may be generated in the early plasma due to anisotropic parton distributions [38, 39]. These large fields, transverse to the beam, tend to deflect radiated gluons from a transversely traveling jet, preferentially,

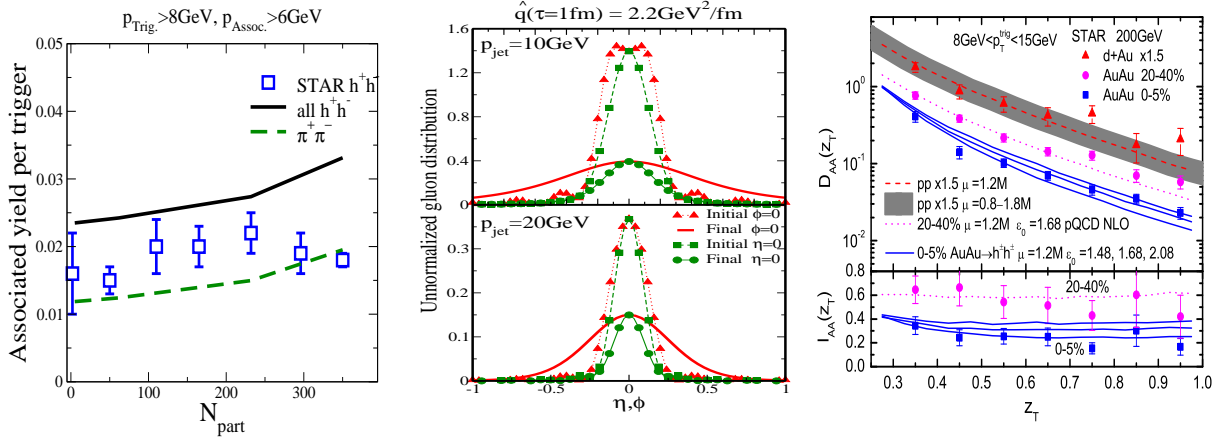


Figure 5. Left: The near side associated yield as a function of centrality of collision [11]. Center: the associated gluon distribution in azimuth and pseudo-rapidity from a hard jet just after production (Initial) and at exit from a medium with turbulent transverse fields (Final) [37]. Right: the triggered associated fragmentation function $D(z_T, p_T^{trig})$ and its ratio in Au-Au with that in $p-p$ [41].

in the longitudinal directions. While such effects influence the solid angle distributions of the radiated gluons around the originating parton, they do not have a considerable effect on the total energy lost. Such phenomena may yield an explanation for the ridge like structure seen in the near side correlations. In a Ref. [37], a quantitative estimate of this broadening was made and the results were found to decrease with increasing energy of the radiated gluon as shown in the center panel of Fig. 5. Here the angular profile of radiated gluons in azimuth (ϕ) and pseudo-rapidity (η) at production and at exit from an expanding medium is plotted. Results are for two energies of the original jet of 10 and 20 GeV, formed at a depth of 3 fm inside the medium. The radiated gluon always carries a fraction $x = 0.4$ of the jet's energy.

In most jet quenching calculations, the hard cross section for both single inclusive and double inclusive observables has to date been computed at leading order [32, 40]. Recently, such calculations have been extended to next-to-leading order (NLO) [41]. Here, the authors computed both single inclusive as well as triggered back-to-back distributions of leading hadrons. The right most panel of Fig. 5 shows the p_T^{trig} weighted differential cross section to produce an associated hadron with a given p_T^{assoc} , given a trigger hadron [also referred to as the triggered fragmentation function $D(z_T, p_T^{trig})$, where $z_T = p_T^{assoc}/p_T^{trig}$]. Also plotted is the ratio of this cross section in Au-Au with that in $p-p$ referred to as the I_{AA} . Such developments along with emerging computations, of heavy quark energy loss and photon hadron correlations hold promise for the central role of jet-medium interactions in heavy-ion physics.

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